MADPH-07-1500 IZTECH-P-07/05 DESY-07-213

Dirac Neutrino Masses from Generalized Supersymmetry Breaking

Durmuş A. Demir*
Department of Physics, Izmir Institute of Technology, IZTECH, TR35430 Izmir, Turkey
Deutsches Elektronen - Synchrotron, DESY, D-22603 Hamburg, Germany

Lisa L. Everett†
Department of Physics, University of Wisconsin, Madison, WI 53706, USA

Paul Langacker‡
School of Natural Sciences, Institute for Advanced Study, Einstein Drive, Princeton, NJ , 08540, USA

We demonstrate that Dirac neutrino masses in the experimentally preferred range are generated within supersymmetric gauge extensions of the Standard Model with a generalized supersymmetry breaking sector. If the usual superpotential Yukawa couplings are forbidden by the additional gauge symmetry (such as a $U(1)'$), effective Dirac mass terms involving the "wrong Higgs" field can arise either at tree level due to hard supersymmetry breaking fermion Yukawa couplings, or at one-loop due to nonanalytic or "nonholomorphic" soft supersymmetry breaking trilinear scalar couplings. As both of these operators are naturally suppressed in generic models of supersymmetry breaking, the resulting neutrino masses are naturally in the sub-eV range. The neutrino magnetic and electric dipole moments resulting from the radiative mechanism also vanish at one-loop order.

arXiv:0712.1341v2 [hep-ph] 19 Dec 2007

The discovery of neutrino oscillations has confirmed that neutrinos are massive and that leptons exhibit nontrivial mixing, providing the first particle physics evidence for physics beyond the Standard Model (SM). Neutrino masses require either the existence of novel matter species not found in the SM spectrum and/or the violation of the global symmetries of the SM via higher-dimensional operators. Extensions incorporating such additional structure should ideally be capable of improving the ultraviolet behavior of the SM beyond Fermi energies. Low-energy softly-broken supersymmetry thus provides a well-motivated theoretical framework in which to incorporate neutrino mass generation mechanisms. As no conclusive experimental indications for neutrinoless double beta decay (for Majorana neutrinos) or neutrino magnetic or electric dipole moments (for Dirac neutrinos) are available at present, we must explore all mechanisms for generating light neutrino masses, not only to reveal the origin of neutrino masses and mixings, but also to determine viable patterns for physics beyond the SM.

Many mechanisms are known for generating light Majorana or Dirac neutrino masses (see e.g. [1, 2, 3, 4]). Some scenarios, including the familiar seesaw mechanism [1], rely upon the supposition that the right-handed neutrinos have no gauge quantum numbers with respect to the low energy gauge group. Right-handed neutrinos are SM gauge singlets, but they can be charged under additional gauge symmetries which may survive from many high-scale theories, such as four-dimensional string models. Thus, if the right-handed neutrinos are not complete (low scale) gauge singlets, these scenarios are not viable, at least not in their simplest implementation. Dirac neutrinos, which occur if lepton number is an exact symmetry, do not necessarily have this requirement. However, as neutrino Dirac masses originate from Yukawa interactions after electroweak breaking, their Yukawa couplings must be exceedingly small. This can be explained if they are forbidden at the renormalizable level by additional symmetries but are generated from higher-dimensional operators. Previous work along these lines (see e.g. [5]) assumes that such operators occur in the superpotential.

In this Letter, we demonstrate that appropriately suppressed Dirac neutrino masses can be generated by generalized supersymmetry breaking terms in models in which the right-handed neutrinos are charged under additional gauge symmetries. These symmetries forbid the usual neutrino superpotential Yukawa couplings, but allow for higher-dimensional operators which lead to effective Dirac neutrino mass terms involving the "wrong Higgs" field upon supersymmetry breaking.

Fermion masses represent the breakdown of chiral flavor symmetries, and as such can be parametrized by vacuum expectation values (VEVs) of scalar fields charged under the flavor symmetry. In theories with low energy supersymmetry, it has long been known [6] (see also [7]) that such chiral flavor symmetries may be broken by the VEVs of auxiliary fields, rather than their scalar counterparts. If the renormalizable superpotential Yukawa couplings and right-handed neutrino Majorana mass terms are forbidden by the symmetry, fermion masses are generated either (i) at tree level due to hard supersymmetry breaking effective Yukawa terms, with

$$m_f \sim Y_{\text{eff}} \langle H \rangle, \tag{1}$$

or (ii) radiatively via sfermion–neutralino loops:

$$m_f \sim \frac{\alpha}{2\pi} \frac{\tilde{A} M_\lambda \langle H \rangle}{\tilde{m}^2}, \tag{2}$$

*email: durmus.demir@desy.de
†email: leverett@wisc.edu
‡email: pgl@ias.edu

in which α denotes a typical gauge coupling, $\tilde{A}$ denotes a soft trilinear scalar coupling, M_λ denotes a gaugino mass, and $\tilde{m}$ denotes a typical sfermion mass.

In generic supersymmetry breaking models, Eq. (1) and Eq. (2) are naturally in the experimentally favored ranges for neutrino masses. The effective Yukawa interaction of Eq. (1) is due to a higher-dimensional Kähler potential operator suppressed by a high scale M (the messenger scale). Hence, $Y_{\rm eff} \sim \tilde{m}/M$, and

$$\left(\frac{m_\nu}{10^{-3}\,\mathrm{eV}}\right) \sim \left(\frac{\tilde{m}}{100\,\mathrm{GeV}}\right)\left(\frac{M}{10^{16}\,\mathrm{GeV}}\right)^{-1}. \qquad (3)$$

Due to the large suppression factor, these effective "wrong-Higgs" Yukawa coupling terms do not spoil the resolution to the hierarchy problem, although technically they are hard supersymmetry breaking operators [8].

Let us now focus on the radiative mass terms of Eq. (2). These terms are suppressed due to the specific trilinear scalar couplings (the $\tilde{A}$ terms) allowed by the flavor symmetry. To understand this suppression, recall that there are two classes of $\tilde{A}$ terms: (i) the standard analytic or "holomorphic" terms, which are coefficients of operators of the form $\phi\phi\phi$, and (ii) the nonanalytic or "nonholomorphic" terms, which accompany $\phi^*\phi\phi$ operators.

In typical models, the nonanalytic trilinear scalar terms, which have previously been considered in the context of radiative SM fermion mass generation [6], are well known to be suppressed by $\tilde{m}/M$ [8, 9]. Recently, it has been claimed that without this strong suppression, Goldstino loops can reintroduce the hierarchy problem [10]. If these terms are so strongly suppressed, they are irrelevant for most phenomenological analyses and cannot provide the dominant contribution to charged fermion masses. However, this suppression is of the right order to be relevant for Dirac neutrino masses:

$$\left(\frac{m_\nu}{10^{-3}\,\mathrm{eV}}\right) \sim \frac{\alpha}{2\pi}\left(\frac{\tilde{m}}{100\,\mathrm{GeV}}\right)\left(\frac{M}{10^{16}\,\mathrm{GeV}}\right)^{-1}, \qquad (4)$$

which can fall within the experimentally allowed range without excessive tuning. Furthermore, the associated radiative neutrino magnetic and electric dipole moments vanish at one loop level.

The nonanalytic terms contribute to quadratic divergences through tadpole diagrams [11], and thus are not soft in the presence of gauge singlets. If SM singlets such as right-handed neutrinos are present these terms can be rendered soft only if the SM gauge group is extended, and all SM singlets are charged under the additional gauge group(s). The simplest extension is to include an additional Abelian $U(1)'$ factor, which can also provide a resolution of the supersymmetric μ problem [12]. The $U(1)'$ charges can be assigned such that the neutrino superpotential Yukawa couplings and the associated trilinear Kähler potential couplings are forbidden, while the wrong-Higgs trilinear couplings are allowed. The nontrivial $U(1)'$ charges of the right-handed neutrinos also forbid bare Majorana mass terms.

We will now provide a detailed analysis of these points. Consider the MSSM augmented by three right-handed neutrino superfields, $\widehat{N}^i = (\widetilde{\nu}_R^{c\,i}, \nu_R^{c\,i})$. Supersymmetry breaking occurs in a hidden sector via the F-component VEV of a chiral superfield $\widehat{X}$, with $\langle\widehat{X}\rangle = F\theta\theta$, and is communicated to the visible sector at a large scale M via nonrenormalizable interactions. The F component of the neutrino superpotential Yukawa coupling then gives an analytic scalar trilinear coupling:[1]

$$\frac{1}{M}\left(\widehat{X}\widehat{L}\cdot\widehat{H}_u\mathbf{Y}_\nu\widehat{N}\right)_F = \widetilde{L}\cdot H_u\mathbf{A}_\nu\widetilde{\nu}_R^c, \qquad (5)$$

with

$$\mathbf{A}_\nu \equiv \frac{F}{M}\mathbf{Y}_\nu \sim \tilde{m}\mathbf{Y}_\nu, \qquad (6)$$

in which $F/M \sim \tilde{m}$ sets the scale of soft-breaking masses (where $\tilde{m} \sim$ TeV). There are also D term contributions from the Kähler potential, which are intrinsically nonanalytic. These contributions lead to suppressed effective Yukawa couplings

$$\frac{1}{M^2}\left(\widehat{X}^\dagger\widehat{L}\cdot\widehat{H}_u\bar{\mathbf{Y}}_\nu\widehat{N}\right)_D = L\cdot H_u\widetilde{\mathbf{Y}}_\nu\nu_R^c, \qquad (7)$$

in which the effective Yukawa coupling is

$$\widetilde{\mathbf{Y}}_\nu \equiv \frac{F}{M^2}\bar{\mathbf{Y}}_\nu \sim \frac{\tilde{m}}{M}\bar{\mathbf{Y}}_\nu, \qquad (8)$$

which were previously studied [13]. They also lead to hard supersymmetry breaking effective fermion Yukawa couplings of the wrong-Higgs form:

$$\frac{1}{M^2}\left(\widehat{X}^\dagger\widehat{L}\cdot\widehat{H}_d^c\bar{\mathbf{Y}}'_\nu\widehat{N}\right)_D = L\cdot H_d^c\widetilde{\mathbf{Y}}'_\nu\nu_R^c, \qquad (9)$$

in which the effective Yukawa coupling is

$$\widetilde{\mathbf{Y}}'_\nu \equiv \frac{F}{M^2}\bar{\mathbf{Y}}'_\nu \sim \frac{\tilde{m}}{M}\bar{\mathbf{Y}}'_\nu. \qquad (10)$$

In addition to the usual scalar mass-squared terms,

$$\frac{1}{M^2}\left(\widehat{X}\widehat{X}^\dagger\widehat{N}^c\mathbf{K}_\nu\widehat{N}\right)_D = \widetilde{\nu}_R^T\mathbf{m}^2_{\widetilde{N}}\widetilde{\nu}_R^c, \qquad (11)$$

with

$$\mathbf{m}^2_{\widetilde{N}} \equiv \frac{F^2}{M^2}\mathbf{K}_\nu \sim \tilde{m}^2\mathbf{K}_\nu, \qquad (12)$$

D terms also lead to wrong-Higgs nonanalytic trilinear couplings:[2]

$$\frac{1}{M^3}\left(\widehat{X}\widehat{X}^\dagger\widehat{L}\cdot\widehat{H}_d^c\mathbf{Y}'_\nu\widehat{N}\right)_D = \widetilde{L}\cdot H_d^c\mathbf{A}'_\nu\widetilde{\nu}_R^c, \qquad (13)$$

[1] The extension to quarks and charged leptons is straightforward.
[2] Unlike the holomorphic couplings, the nonholomorphic couplings are independent of the superpotential.

with

$$\mathbf{A}'_\nu \equiv \frac{F^2}{M^3}\mathbf{Y}'_\nu \sim \frac{\tilde{m}^2}{M}\mathbf{Y}'_\nu. \tag{14}$$

Hence, $\mathbf{A}'_\nu$ is suppressed by $F/M^2 = \tilde{m}/M$ with respect to $\mathbf{A}_\nu \sim F/M$. It is the F/M^2 suppression which plays a key role in neutrino mass generation in both cases. The F/M^2 suppression has been discussed previously [13, 14]; however, these works present models in which nonholomorphic terms lead to Majorana masses and holomorphic operators lead to Dirac masses, and do not typically allow for the right-handed neutrinos to have nontrivial charges under additional gauge symmetries.

To allow the wrong-Higgs couplings of Eq. (9) and Eq. (13) and forbid the usual neutrino Yukawa couplings (both tree level and effective, as in Eq. (5) and Eq.(7)), we assume that the right-handed neutrinos are charged under an extended gauge group. This prevents $\widehat{N}^i$ from acquiring a large tree-level Majorana mass[3] (in contrast to the seesaw mechanism), and has the added advantage that the nonanalytic trilinear couplings of Eq. (13) now are soft breaking terms (*i.e.*, no quadratic divergences are induced in the scalar sector). The simplest gauging, though not the only logical possibility, is to add a new Abelian gauge factor $U(1)'$, with charges that satisfy

$$\begin{aligned} Q_L + Q_{H_u} + Q_N &\neq 0, && (15)\\ Q_L - Q_{H_d} + Q_N &= 0. && (16)\end{aligned}$$

These conditions are clearly inconsistent with having a bare superpotential μ term. The remedy is to replace the μ parameter by a chiral SM singlet $\widehat{S}$ with a nonvanishing $U(1)'$ charge Q_S, with $Q_S + Q_{H_u} + Q_{H_d} = 0$ [12], such that an effective μ term is induced by the VEV of S.[4] In this case, it is worth noting that upon $U(1)'$ breaking, superpotential holomorphic couplings of the form

$$\frac{1}{M}\widehat{S}\widehat{L}\cdot\widehat{H}_u\mathbf{Y}''_\nu\widehat{N} \tag{17}$$

may also be generated. As discussed in [5], these may give rise to an additional ("right-Higgs") contribution to the Dirac masses of a similar order of magnitude:

$$m_f = \frac{\langle S\rangle}{M}\mathbf{Y}''_\nu\langle H_u^0\rangle \sim \frac{\tilde{m}}{M}\mathbf{Y}''_\nu\langle H_u^0\rangle. \tag{18}$$

We assume any $U(1)'$ gauge anomalies are cancelled by GUT remnants at the TeV scale; one can also consider anomaly free family-dependent $U(1)'$ groups [7].

We now turn to a more precise analysis of the neutrino masses generated by Eq. (9) and Eq. (13). The Yukawa interaction Eq. (9) induces a Dirac neutrino mass

$$m_\nu = \langle H_d^0\rangle\widetilde{\mathbf{Y}}'_\nu, \tag{19}$$

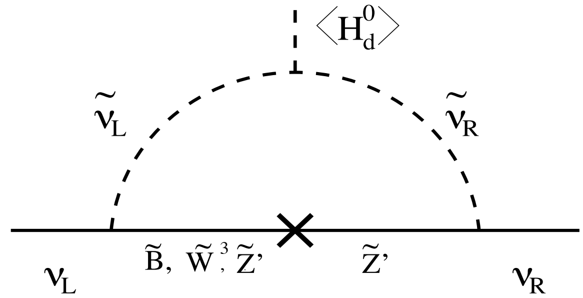

FIG. 1: The one-loop diagram that generates radiative Dirac neutrino masses.

in agreement with Eq. (1) and Eq. (3). This interaction is technically hard, but the resulting Higgs mass shift $\delta m^2_{H_d} = -(1/(8\pi^2))\widetilde{\mathbf{Y}}'^\dagger_\nu\widetilde{\mathbf{Y}}'_\nu M^2 = -(1/(8\pi^2))\tilde{m}^2\bar{\mathbf{Y}}'^\dagger_\nu\bar{\mathbf{Y}}'_\nu$ is too small to leave any impact on the gauge hierarchy.

For the radiatively induced neutrino masses, the requisite Lagrangian terms are

$$\frac{g_Y}{\sqrt{2}}\tilde{\nu}_L^\dagger\mathcal{N}_i\chi_i^0\nu_L + \sqrt{2}g'_Y Q_N\tilde{\nu}_R^T N^0_{Z'i}\chi_i^0\nu_R^c + \text{h.c.}, \tag{20}$$

in which $N^0_{\eta^0 i}$ denotes the contamination of the neutralino gauge eigenstate $\eta^0 \in \{\widetilde{Z}', \widetilde{B}, \widetilde{W}^3, \widetilde{H}_d^0, \widetilde{H}_u^0, \widetilde{S}\}$ in the ith neutralino χ_i^0 $(i = 1,\ldots,6)$, and $\mathcal{N}_i$ is

$$\mathcal{N}_i = \cot\theta_W N^0_{W^3 i} - N^0_{Bi} + 2Q_L\frac{g'_Y}{g_Y}N^0_{Z'i}. \tag{21}$$

These interactions induce Dirac neutrino masses at one loop, as shown in Figure 1, of the form:[5]

$$\begin{aligned}\mathbf{m}_{\nu\,ab} &= \frac{g_Y g'_Y\langle H_d^0\rangle Q_N}{32\pi^2}\Big\{\mathcal{S}_{L\,ac}(\mathcal{S}_L^\dagger\mathbf{A}'_\nu\mathcal{S}_R)_{cd}\mathcal{S}^\dagger_{R\,db}\\ &\quad\times\ m_{\chi_i^0}N^0_{Z'i}\mathcal{N}_i F\left(m^2_{\tilde{\nu}_{L\,c}}, m^2_{\tilde{\nu}_{R\,d}}, m^2_{\chi_i^0}\right)\Big\}, \end{aligned}\tag{22}$$

in which repeated indices are summed over, and $\mathcal{S}_L$ and $\mathcal{S}_R$ are the sneutrino mixing matrices,[6] defined via

$$\begin{aligned}\mathcal{S}_L^\dagger\mathbf{m}^2_{\tilde{\nu}_L}\mathcal{S}_L &= \text{diag.}\left(m^2_{\tilde{\nu}_{L\,1}}, m^2_{\tilde{\nu}_{L\,2}}, m^2_{\tilde{\nu}_{L\,3}}\right) && (24)\\ \mathcal{S}_R^T\mathbf{m}^2_{\tilde{\nu}_R}\mathcal{S}_R^* &= \text{diag.}\left(m^2_{\tilde{\nu}_{R\,1}}, m^2_{\tilde{\nu}_{R\,2}}, m^2_{\tilde{\nu}_{R\,3}}\right). && (25)\end{aligned}$$

---

[3] See [15] for related work involving discrete gauge symmetries.

[4] One can also require that charged fermion masses are generated radiatively, which requires much larger soft trilinear couplings.

[5] Due to small mixing, the $\widetilde{B}$ and $\widetilde{W}^3$ contributions are typically subdominant to that of the $\widetilde{Z}'$.

[6] Their mass-squares are obtained by adding the associated D-term contributions

$$\begin{aligned}\mathbf{m}^2_{\tilde{\nu}_L} &= \mathbf{m}^2_{\tilde{L}} + \frac{1}{2}\cos 2\beta M_Z^2 + \frac{1}{2}Q_L\delta^2_{Z'}\\ \mathbf{m}^2_{\tilde{\nu}_R} &= \mathbf{m}^2_{\tilde{N}} + \frac{1}{2}Q_N\delta^2_{Z'}, \end{aligned}\tag{23}$$

with $\delta^2_{Z'} = 2{g'_Y}^2\left(Q_{H_u}\langle H_u^0\rangle^2 + Q_{H_d}\langle H_d^0\rangle^2 + Q_S\langle S\rangle^2\right)$. $\langle S\rangle$ sets the effective μ parameter below the $U(1)'$ breaking scale [12].

The loop function appearing in Eq.(22) is given by

$$F\left(m_1^2, m_2^2, m^2\right) = \frac{1}{m_1^2 - m_2^2}\left(\frac{\ln\beta_1}{\beta_1 - 1} - \frac{\ln\beta_2}{\beta_2 - 1}\right). \quad (26)$$

$\beta_i = m^2/m_i^2$ reduces to $1/2m^2$ when $m_1 = m_2 = m$.

Eq. (22) demonstrates that neutrinos acquire Dirac masses radiatively only if the right-handed neutrinos are gauged under the $U(1)'$ symmetry. $U(1)'$ invariance thus not only ensures that the nonanalytic trilinear terms are soft, but also provides the the chirality flip required for neutrino mass generation through the $\widetilde{Z}'$, which couples to both left- and right-handed neutrinos.

For $M \sim M_{GUT}$, the neutrino masses are in the right range (the $\alpha/2\pi$ suppression can be countered by relaxing the degeneracy among the superpartner masses; this factor is absent for the tree-level masses of Eq. (19)). If $M \sim M_{Pl}$, an enhancement is required. For other mediation mechanisms the messenger scale can be lowered, depending on the details of the model.

The flavor structure of the tree-level Dirac neutrino mass Eq. (19) depends only on $\bar{\mathbf{Y}}'_\nu$ in Eq.(9). However, the flavor structure of the radiative neutrino masses involves $\mathbf{m}^2_{\tilde{\nu}_L}$, $\mathbf{m}^2_{\tilde{\nu}_R}$, and $\mathbf{A}'_\nu$. If the left-handed and right-handed sneutrinos are approximately degenerate in mass, the neutrino mixings are controlled by the nonanalytic trilinear coupling $\mathbf{A}'_\nu$ alone. Alternatively, $\mathbf{A}'_\nu$ may be strictly diagonal, such that neutrino mixings arise from nontrivial flavor structures of $\mathbf{m}^2_{\tilde{\nu}_L}$ and $\mathbf{m}^2_{\tilde{\nu}_R}$.

The radiative mechanism that leads to fermion masses also generically induces electric and magnetic dipole moments [6, 16]. However, in this scenario, the neutrino dipole moments vanish at one loop. This occurs because the right-handed neutrinos do not couple directly to the higgsinos through Yukawa interactions, and they do not have any charged gaugino with which to interact. Dirac neutrino masses also induce dipole moments within the SM of order $10^{-19}\mu_B$, which are much smaller than the best available bounds (of order $10^{-12}\mu_B$) [17].

In this Letter, we have discussed mechanisms to induce naturally suppressed neutrino Dirac masses within gauge-extended models with low energy supersymmetry. Neutrino mass terms of the "wrong-Higgs" type are generated either at tree level from formally hard (but in practice safe) effective Yukawa couplings, or radiatively due to nonanalytic soft supersymmetry breaking interactions. The neutrino mass scale naturally falls within the experimentally allowed range due to the $F/M^2 \sim \tilde{m}/M$ suppression. Moreover, this mechanism is operational for models in which the right-handed neutrinos are not complete singlets of the low energy gauge group. This scenario, apart from providing an understanding of the origin of naturally suppressed Dirac neutrino masses, allows for a natural resolution of the supersymmetric μ problem and leads to TeV-scale $U(1)'$ physics which should be testable at forthcoming colliders such as the LHC.

Acknowledgments

D.D. is supported by Alexander von Humboldt-Stiftung Friedrich Wilhelm Bessel-Forschungspreise and by the Turkish Academy of Sciences via the GEBIP grant. L.E. is supported by the U.S. Department of Energy grant DE-FG02-95ER40896. P.L. is supported by Friends of the IAS and by the NSF grant PHY-0503584. We thank D. J. Chung, A. Ibarra, I. W. Kim, C. Luhn, N. Seiberg, and L. T. Wang for fruitful discussions.

[1] P. Minkowski, Phys. Lett. B **67**, 421 (1977); T. Yanagida, in *Proceedings of the Workshop on the Unified Theory and the Baryon Number in the Universe* (O. Sawada and A. Sugamoto, eds.), KEK, Tsukuba, Japan, 1979, p. 95; M. Gell-Mann, P. Ramond, and R. Slansky, *Complex spinors and unified theories*, in *Supergravity* (P. van Nieuwenhuizen and D. Z. Freedman, eds.), North Holland, Amsterdam, 1979, p. 315.

[2] S. Weinberg, Phys. Rev. Lett. **43**, 1566 (1979).

[3] A. Y. Smirnov, arXiv:hep-ph/0411194.

[4] See e.g. the review R. N. Mohapatra *et al.*, Rept. Prog. Phys. **70**, 1757 (2007) [arXiv:hep-ph/0510213].

[5] G. Cleaver, M. Cvetic, J. R. Espinosa, L. L. Everett and P. Langacker, Phys. Rev. D **57**, 2701 (1998) [arXiv:hep-ph/9705391]; P. Langacker, Phys. Rev. D **58**, 093017 (1998) [arXiv:hep-ph/9805281].

[6] F. Borzumati, G. Farrar, N. Polonsky and S. Thomas, Nucl. Phys. B **555**, 53 (1999) [arXiv:hep-ph/9902443].

[7] D. A. Demir, G. L. Kane and T. T. Wang, Phys. Rev. D **72**, 015012 (2005) [arXiv:hep-ph/0503290].

[8] M. Dine, N. Seiberg and S. Thomas, Phys. Rev. D **76**, 095004 (2007) [arXiv:0707.0005 [hep-ph]].

[9] L. Hall and L. Randall, Phys. Rev. Lett. **65**, 2939 (1990); I. Jack and D. R. T. Jones, Phys. Lett. B **457**, 101 (1999) [arXiv:hep-ph/9903365]; S. Martin, Phys. Rev. D **61**, 035004 (2000) [arXiv:hep-ph/9907550]; N. Polonsky and S. f. Su, Phys. Rev. D **63**, 035007 (2001) [arXiv:hep-ph/0006174]; P. Camara, L. Ibanez and A. Uranga, Nucl. Phys. B **689**, 195 (2004) [arXiv:hep-th/0311241].

[10] J. A. Bagger and A. F. Falk, arXiv:0708.3364 [hep-ph].

[11] L. Girardello, M. Grisaru, Nucl. Phys. B **194**, 65 (1982).

[12] D. Suematsu and Y. Yamagishi, Int. J. Mod. Phys. A **10**, 4521 (1995) [arXiv:hep-ph/9411239]; M. Cvetic, D. A. Demir, J. R. Espinosa, L. L. Everett and P. Langacker, Phys. Rev. D **56**, 2861 (1997) [Erratum-ibid. D **58**, 119905 (1998)] [arXiv:hep-ph/9703317].

[13] N. Arkani-Hamed, L. J. Hall, H. Murayama, D. R. Smith and N. Weiner, Phys. Rev. D **64**, 115011 (2001) [arXiv:hep-ph/0006312]; F. Borzumati and Y. Nomura, Phys. Rev. D **64**, 053005 (2001) [arXiv:hep-ph/0007018]; F. Borzumati, K. Hamaguchi, Y. Nomura and T. Yanagida, arXiv:hep-ph/0012118; J. March-Russell and S. West, Phys. Lett. B **593**, 181 (2004) [arXiv:hep-ph/0403067].

[14] G. Dvali and Y. Nir, JHEP **9810**, 014 (1998) [arXiv:hep-ph/9810257]; J. Frere, M. Libanov and S. Troitsky, Phys. Lett. B **479**, 343 (2000) [arXiv:hep-ph/9912204]; J. Casas, J. R. Espinosa and I. Navarro, Phys. Rev. Lett. **89**, 161801 (2002) [arXiv:hep-ph/0206276].

[15] C. Luhn and M. Thormeier, arXiv:0711.0756 [hep-ph].

[16] K. Fujikawa and R. Shrock, Phys. Rev. Lett. **45**, 963 (1980).

[17] A. Kuznetsov and N. Mikheev, arXiv:0709.0110 [hep-ph].